NMSSM with Lopsided Gauge Mediation

Ivan Donkin, Alexander K. Knochel

Institut für Theoretische Physik, Universität Heidelberg, Philosophenweg 19,
D-69120 Heidelberg, Germany

(donkin@thphys.uni-heidelberg.de a.k.knochel@thphys.uni-heidelberg.de)

Abstract

We study a gauge mediated supersymmetry breaking version of the NMSSM in which the soft $m_{H_u}^2$ and $m_{H_d}^2$ masses receive extra contributions due to the presence of direct couplings between the Higgs and the messenger sector. We are motivated by the well-known result that minimal gauge mediation is phenomenologically incompatible with the NMSSM due to the small value of the induced effective μ term. The model considered in the present paper solves the aforementioned problem through a modified RG running of the singlet soft mass m_N^2. This effect, which is induced by the dominant $m_{H_d}^2$ term in the one-loop β-function of m_N^2, shifts the singlet soft mass towards large negative values at the electroweak scale. That is sufficient to ensure a large VEV for the scalar component of the singlet which in turn translates into a sizeable effective μ term. We also describe a mechanism for generating large soft trilinear terms at the messenger scale. This allows us to make the mass of the lightest Higgs boson compatible with the current LHC bound without relying on exceedingly heavy stops.

1 Introduction

Gauge mediation constitutes a very predictive framework which describes the transmission of SUSY breaking effects from the hidden to the observable sector [1–3]. Among other things it guarantees flavor universality of the soft sfermion masses at the messenger scale and therefore suppresses the potentially dangerous flavor-changing neutral currents [4].

However, in its minimal form gauge mediation has nothing to say about the origin of the μ and B_μ terms. The problem arises as follows: Recall that in MGM models SUSY breaking effects are transmitted from the hidden to the observable sector only through gauge interactions. In the absence of direct couplings between the messenger and the Higgs sector, Peccei-Quinn is an exact symmetry of the fundamental theory. The Higgs bilinear term $\mu H_u H_d$, on the other hand, breaks this symmetry explicitly [5,6]. Therefore this operator cannot be present in the low-energy effective action which arises after integrating out the heavy messenger fields at one loop. In this sense the MGM scenario is not considered satisfactory unless one specifies the extra mechanism responsible for generating μ and B_μ.

There are various proposals in the literature for generating a Higgs bilinear term of the correct order of magnitude. In the context of the so called Giudice-Masiero solution, which was initially formulated in models with gravity mediated supersymmetry breaking (see [7]), a tree-level μ-term is forbidden from the original superpotential and is induced by the following higher-dimensional effective operators in the Kähler potential:

$$K \supset \frac{c_1}{M} \int d^4\theta H_u H_d X^\dagger \;+\; \frac{c_2}{M^2} \int d^4\theta H_u H_d X X^\dagger \;+\; \text{h.c.} \tag{1}$$

The parametrization in eq.(1) applies both to models with gauge and gravity mediated supersymmetry breaking. In the GMSB context X is a spurion superfield which couples to the hidden sector, M is the messenger scale and the prefactors c_1 and c_2 stand for the product of coupling constants and possible loop factors. The interactions in eq.(1) can be generated radiatively after introducing direct couplings between the Higgs and the messenger sector and integrating out the heavy messenger superfields at one loop. Substituting X by its F-term VEV, we obtain a $\mu = c_1 \Lambda$ term from the first operator in eq.(1) and a $B_\mu = c_2 \Lambda^2$ term from the second operator (with $\Lambda = F_X/M$ being the effective scale of SUSY breaking). In theories with gravity mediated supersymmetry breaking one can assume that the prefactors c_1 and c_2 are of order one, $c_1 \sim c_2 \sim \mathcal{O}(1)$, which leads to the celebrated Guidice-Masiero solution $B_\mu \sim \mu^2$. In the gauge mediated supersymmetry breaking setup, on the other hand, μ and B_μ are generated at one loop and one anticipates that $c_1 \sim c_2 \sim 1/16\pi^2$. Therefore the second power μ^2 is loop suppressed with respect to B_μ, i.e. $B_\mu \sim 16\pi^2 \mu^2$. This leads to the relation

$$B_\mu \gg \mu^2 \tag{2}$$

which is incompatible with electroweak symmetry breaking due to the unacceptably large B_μ. For this reason eq.(2) is usually referred to as the μ/B_μ problem of gauge-mediated supersymmetry breaking [8].

Alternative solutions, which avoid eq.(2) and ensure that $B_\mu \sim \mu^2$, were formulated e.g. in the context of the so called dynamical relaxation mechanism [8]. More recently, the authors of [9] have pointed out that successful EWSB can be achieved even if the problematic relation $B_\mu \gg \mu^2$ is kept intact provided that one assumes a non-trivial hierarchy between the mass terms in the Higgs sector at the electroweak scale:

$$\mu^2 \sim m_{H_u}^2 \ll B_\mu \ll m_{H_d}^2 \ . \tag{3}$$

As was argued in [9] the pattern in eq.(3) can be obtained by considering models with gauge mediated supersymmetry breaking in which the two Higgs doublets couple directly to messenger superfields (or more generally to superfields in the SUSY breaking sector).

In the present paper we analyze eq.(3) in the context of the Next to Minimal Supersymmetric Standard Model (NMSSM) [29]. We are motivated by the well-known result that the NMSSM is incompatible with the minimal form of gauge mediated supersymmetry breaking [32] due to the small value of the generated μ term. We will argue that a phenomenologically viable model can be obtained by a slight deformation of the hierarchy between the soft mass terms in the Higgs sector. To be exact, we consider an NMSSM in which the $m_{H_d}^2$ term dominates all other soft masses, in particular we require that $m_{H_d}^2 \gg m_{H_u}^2$. We describe an embedding of our scenario into a GMSB setup with direct Higgs-messenger mixing and refer to this model as lopsided NMSSM in accord with the terminology introduced in [10]. The crucial difference between our approach and the more conventional lopsided MSSM is the origin of the effective μ term. In the minimal setup considered in [10] μ arises solely from the tree-level Higgs-messenger couplings. In our setting there is a large extra contribution coming from the NMSSM operator $\lambda N H_u H_d$. As a result the μ term in the lopsided NMSSM is far less constrained than its MSSM counterpart.

The paper is organized as follows: In section 2 we introduce the lopsided GMSB scenario following closely the approach in [9]. Using this line of reasoning we define a lopsided NMSSM which is the main object of study in the remainder of the paper. In section 3 we give a detailed description of the field content and superpotential couplings of our theory, including a full specification of the hidden and messenger sectors. In section 4 we show that the effective μ term is comprised of two pieces – one arising from the Higgs-messenger mixing and another associated with the NMSSM part of the superpotential. We describe in detail the mechanism which triggers a large vacuum expectation value for the NMSSM singlet and therefore generates the latter piece of the Higgsino mass parameter. In section 5 we calculate the low-energy spectrum for several points in parameter space. Crucially, we describe a new way for generating large soft trilinear terms at the messenger scale. As pointed out already in [31] this is chiefly important for raising the Higgs mass above the current LHC bound. We also find that, for large regions of parameter space, the negative tree-level contribution to m_{h^0} from the mixing with the singlet is unacceptably large and discuss a possible mechanism which can suppress this effect. In section 6 we briefly discuss the collider phenomenology of the model.

2 Defining the lopsided NMSSM

In order to set the stage for our subsequent discussion and fix the notation we briefly recall the basics of lopsided gauge mediation. Here we follow the approach in [9], i.e. we introduce the lopsided MSSM as a possible solution to the μ/B_μ problem. Let us analyze eq.(3) in more detail. To this end we shall look at the tree-level minimization conditions in the Higgs sector:

$$\frac{m_Z^2}{2} = -|\mu|^2 - \frac{m_{H_u}^2 \tan^2 \beta - m_{H_d}^2}{\tan^2 \beta - 1} \tag{4}$$

$$\sin 2\beta = \frac{2B_\mu}{2|\mu|^2 + m_{H_u}^2 + m_{H_d}^2} \tag{5}$$

One can argue that the above hierarchy leads to a fully natural electroweak symmetry breaking. Indeed, taking into consideration that $\sin 2\beta \approx 2B_\mu/m_{H_d}^2$, we can rewrite (5) as a quadratic equation for $\tan \beta$ which leads to

$$\tan \beta \approx \frac{m_{H_d}^2}{B_\mu} \gg 1 \,. \tag{6}$$

This shows that the proposed solution is operational only in the large $\tan \beta$ regime. From eq.(4) we then obtain:

$$\frac{m_Z^2}{2} = -|\mu|^2 - m_{H_u}^2 - \frac{m_{H_d}^2}{\tan^2 \beta} \tag{7}$$

Note that one can make all terms on the right hand side roughly of the same order of magnitude, while still keeping the hierarchy in eq.(3) intact. This is most easily seen by considering the parametrization from [9]

$$\mu = \epsilon \Lambda_{\text{Higgs}} \qquad B_\mu = \epsilon \Lambda_{\text{Higgs}}^2 \qquad m_{H_u}^2 = \epsilon^2 \Lambda_{\text{Higgs}}^2 \qquad m_{H_d}^2 = \Lambda_{\text{Higgs}}^2 \tag{8}$$

where $\epsilon \ll 1$ is some small number and Λ_{Higgs} denotes the effective scale of SUSY breaking in the Higgs sector. Noting that $\tan \beta = 1/\epsilon$ we deduce that all terms on the right hand side of eq.(7) are of the order $\epsilon^2 \Lambda_{\text{Higgs}}^2$. Thus we expect that in the absence of large radiative correction electroweak symmetry breaking should occur naturally, without the need for significant fine-tuning.

The authors of [9] provide a possible mechanism for generating the pattern in eq.(3). Specifically, we can envisage a non-minimal gauge-mediated scenario with direct couplings between the Higgs and the messenger sector of the form

$$W \supset \lambda_u H_u \Phi_1 \bar{\Phi}_2 + \lambda_d H_d \bar{\Phi}_1 \Phi_2 + \text{h.c} \,. \tag{9}$$

Here $\Phi_1 \bar{\Phi}_2$ and $\bar{\Phi}_1 \Phi_2$ are bilinears of messenger fields. After integrating out the the heavy fields represented by the vector-like pairs Φ_1, $\bar{\Phi}_1$ and Φ_2, $\bar{\Phi}_2$ we obtain the following contributions to the mass parameters in the Higgs sector

$$\mu \approx \lambda_u \lambda_d \frac{N_{\text{Higgs}}}{16\pi^2} \Lambda_{\text{Higgs}} \,, \quad B_\mu \approx \lambda_u \lambda_d \frac{N_{\text{Higgs}}}{16\pi^2} \Lambda_{\text{Higgs}}^2 \,, \quad m_{H_{u,d}}^2 \approx \lambda_{u,d}^2 \frac{N_{\text{Higgs}}}{16\pi^2} \Lambda_{\text{Higgs}}^2 \tag{10}$$

where N_{Higgs} parametrizes the number of messenger fields coupled to the Higgs sector (i.e. $N_{\text{Higgs}} = n_{5+\bar{5}} + 3n_{10+\overline{10}}$ is the "messenger index" in the Higgs sector). Observe that Λ_{Higgs} could but does not have to coincide with the SUSY breaking scale Λ from the MGM sector. The same statement applies to N_{Higgs}. Imposing the hierarchy $\lambda_d \gg \lambda_u$ and introducing the parameter $\epsilon = \frac{N_{\text{Higgs}}}{16\pi^2} \frac{\lambda_d}{\lambda_u}$ leads to the correct low-energy pattern from eq.(3). In [10] the class of models specified by eqs.(3) and (9) was labeled lopsided gauge mediation.

In the present section we introduce the lopsided NMSSM as an effective theory, i.e. we do not care how supersymmetry is broken or transmitted to the observable sector. To begin the description of our model let us fix the superpotential in the Higgs sector:

$$W \supset (\mu + \lambda N)H_u H_d + \frac{\kappa}{3} N^3 \ . \tag{11}$$

Recall that the operator $\lambda N H_u H_d$ is invariant under a Peccei-Quinn symmetry which acts on the N, H_u, H_d superfields according to

$$H_u \to H_u \, e^{-i\phi} \qquad H_d \to H_d \, e^{-i\phi} \qquad N \to N e^{i2\phi} \ . \tag{12}$$

This symmetry is broken at tree-level in a twofold manner – both by the cubic $\frac{\kappa}{3} N^3$ operator and the bilinear mass term for the H_u, H_d Higgs doublets. As a result the potentially dangerous Peccei-Quinn axion in the Higgs sector is automatically avoided. There is an additional subtlety – the $\frac{\kappa}{3} N^3$ term is invariant under a residual $\mathbb{Z}_3$ symmetry (the action of the $\mathbb{Z}_3$ group rotates all three superfields N, H_u and H_d by the same phase $e^{2\pi i/3}$). In the standard $\mathbb{Z}_3$-invariant NMSSM this symmetry is responsible for the appearance of domain walls [19, 21, 29]. Note, however, that in the setup we are considering the discrete $\mathbb{Z}_3$ is not an actual symmetry of the Lagrangian – it is broken at tree-level by the presence of a μ-term. Therefore the domain wall problem is not present in our version of the lopsided NMSSM scenario. [1]

The soft supersymmetry breaking terms in the Higgs sector read

$$-\mathcal{L}_{\text{soft}} = \quad m_N^2|N|^2 + m_{H_u}^2|H_u|^2 + m_{H_d}^2|H_d|^2 + (\, \lambda A_\lambda N H_u H_d \, + \, \frac{1}{3}\,\kappa A_\kappa N^3 \, + \, \text{h.c.}\,) +$$
$$+ \, (\, B_\mu H_u H_d \, + \, \text{h.c.}\,) \tag{13}$$

Imposing $m_{H_d}^2 \gg m_{H_u}^2$ *at the electroweak scale* completes the description of the effective Lagrangian. From a low-energy point of view the precise mechanism which produces the extra contributions to $m_{H_d}^2$ and $m_{H_u}^2$ to ensure the non-trivial hierarchy between the soft mass terms in the Higgs sector is unimportant. The same statement applies to the origin of the μ and B_μ terms in eqs.(11) and (13). We can simply assume that they are present in the Lagrangian and then extract their values from the tree-level minimization conditions in the Higgs sector. In the next section we will construct a high-energy model which produces the lopsided NMSSM as its low-energy limit.

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[1] Clearly one can consider a lopsided NMSSM without a tree-level μ term. In this case the domain wall problem has to be solved through an alternative mechanism.

3 High-energy completion in the context of gauge mediation

In this section we will embed the lopsided NMSSM into a gauge-mediated scenario. Our guiding principle is very simple – the model we are considering should, upon integrating out the messenger superfields at one loop, produce the lopsided NMSSM as its low-energy limit. The main building block of our construction is the following superpotential considered in [10]:

$$W \; \supset \; X\Phi\bar{\Phi} + \lambda N H_d H_u + \frac{\kappa}{3} N^3 + \tag{14}$$

$$+ \; \lambda_u H_u D\,T + \lambda_d H_d \bar{D}\,\bar{T} + X_D D\bar{D} + \frac{X_T}{2}\,(\,T\ \bar{T}\,)\begin{pmatrix} a_T & a_{T\bar{T}} \\ a_{T\bar{T}} & a_{\bar{T}} \end{pmatrix}\begin{pmatrix} T \\ \bar{T} \end{pmatrix}$$

The operator $X\Phi\bar{\Phi}$ comprises the universal part of the messenger sector which is present in any model of gauge mediation. This part consists of messenger superfields $\Phi, \bar{\Phi}$ filling complete vector-like representation of the GUT group $SU(5)$. The masses and splittings between the $\Phi, \bar{\Phi}$ are controlled by a spurion superfield X which acquires VEVs in its scalar and F-term components due to some dynamics in the hidden sector:

$$X = M + F\,\theta^2 \; . \tag{15}$$

After integrating out $\Phi, \bar{\Phi}$ at the messenger scale M, gaugino masses are generated at one and sfermion masses at two loop

$$M_i = \frac{\alpha_i}{4\pi}\,n\,\Lambda \qquad m_{\tilde{f}}^2 = 2\,n\,\Lambda^2\left[\sum_{i=1}^{3} C_i^{\tilde{f}}\left(\frac{\alpha_i}{4\pi}\right)^2\right] \; . \tag{16}$$

Here $\Lambda = F/M$ is the effective SUSY breaking scale and n stands for the messenger index. The $C_i^{\tilde{f}}$ denotes the quadratic Casimir operator for the respective representation of the sfermion superfield $\tilde{f}$ with respect to the SM gauge group G_i, where $G_1 = U(1)$, $G_2 = SU(2)$ and $G_3 = SU(3)$.

The model contains an NMSSM-like sector with an extra singlet N which couples to the Higgs doublets through the usual NMSSM operator $\lambda N H_u H_d$. To break the PQ symmetry in the Higgs sector we introduce the cubic term $\frac{k}{3}N^3$ to the superpotential. At this stage the soft Higgs masses $m_{H_d}^2$ and $m_{H_u}^2$ are equal at the messenger scale. The extra contributions to $m_{H_d}^2$ and $m_{H_u}^2$ which fix the hierarchy $m_{H_d}^2 \gg m_{H_u}^2$ come from direct couplings between the Higgs and the messenger sector $\lambda_u H_u DT + \lambda_d H_d \bar{D}\bar{T}$. Here $D, \bar{D}$ are messenger superfields in the fundamental representation of $SU(2)$ with opposite hypercharges: $D = (\mathbf{1}, \mathbf{2})_1$, $\bar{D} = (\mathbf{1}, \mathbf{2})_{-1}$ (where the subscript stands for the $U(1)_Y$ charge). In addition we have two total singlets T and $\bar{T}$ with quantum numbers $T = (\mathbf{1}, \mathbf{1})_0$ and $\bar{T} = (\mathbf{1}, \mathbf{1})_0$. In order to preserve gauge coupling unification we assume that the $D, \bar{D}$ pair is lying in one $\mathbf{5} \oplus \bar{\mathbf{5}}$ copy of $SU(5)$. As usual we make the assumption that the two spurions acquire VEVs in their scalar and F-term components

$$X_D = M_D(\,1 + \Lambda_D\theta^2) \qquad\qquad X_T = M_T(\,1 + \Lambda_T\theta^2) \; . \tag{17}$$

The messenger sector of the above superpotential (the second line in eq.(14)) exhibits a discrete $\mathbb{Z}_2$ symmetry in the limit $a_{T\bar{T}} = 0$. The action of the $\mathbb{Z}_2$ group is given by:

$$H_u \rightarrow -H_u \qquad T \rightarrow -T \tag{18}$$

Note that this symmetry forbids a μ (and a corresponding B_μ term) in the effective low-energy action. This means that we can control the magnitude of μ and B_μ by suitably adjusting the value of the $a_{T\bar{T}}$ coupling. Following [10] we can diagonalize the $T, \bar{T}$ mass matrix with a suitably chosen rotation by an angle Θ. In the following we shall denote the eigenvalues of the mass matrix by p and ξp where ξ is the proportionality factor. The contribution to the mass parameters in the Higgs sector coming from the superpotential (14) have been calculated in [10]. Explicitly they read

$$m_{H_u}^2 = \frac{\lambda_u^2}{16\pi^2} \Lambda_D^2 \left(\cos^2 \Theta \, P(x,y) + \sin^2 \Theta \, P(\xi x, y) \right) = a_u \frac{\lambda_u^2}{16\pi^2} \Lambda_D^2 \tag{19}$$

$$m_{H_d}^2 = \frac{\lambda_d^2}{16\pi^2} \Lambda_D^2 \left(\sin^2 \Theta \, P(x,y) + \cos^2 \Theta \, P(\xi x, y) \right) = a_d \frac{\lambda_d^2}{16\pi^2} \Lambda_D^2 \tag{20}$$

$$\mu = \frac{\lambda_u \lambda_d}{16\pi^2} \Lambda_D \sin \Theta \cos \Theta \left(-Q(x,y) + Q(\xi x, y) \right) = a_\mu \frac{\lambda_u \lambda_d}{16\pi^2} \Lambda_D \tag{21}$$

$$B_\mu = \frac{\lambda_u \lambda_d}{16\pi^2} \Lambda_D^2 \sin \Theta \cos \Theta \left(-R(x,y) + R(\xi x, y) \right) = a_{B_\mu} \frac{\lambda_u \lambda_d}{16\pi^2} \Lambda_D^2 \tag{22}$$

$$A_u = \frac{\lambda_u^2}{16\pi^2} \Lambda_D \left(\cos^2 \Theta \, S(x,y) + \sin^2 \Theta \, S(\xi x, y) \right) = a_{A_u} \frac{\lambda_u \lambda_d}{16\pi^2} \Lambda_D \tag{23}$$

$$A_d = \frac{\lambda_d^2}{16\pi^2} \Lambda_D \left(\sin^2 \Theta \, S(x,y) + \cos^2 \Theta \, S(\xi x, y) \right) = a_{A_d} \frac{\lambda_u \lambda_d}{16\pi^2} \Lambda_D \tag{24}$$

where we used the shorthand notation $x = M_T/M_D$ and $y = \Lambda_T/\Lambda_D$. The coefficient functions $P(x,y)$, $Q(x,y)$, $R(x,y)$ and $S(x,y)$ appearing in eqs.(19)–(24) are listed in Appendix B. We can recover the $\mathbb{Z}_2$-symmetric limit by setting $a_{T\bar{T}} = 0$. The crucial feature of the model we are considering is that it allows us to switch off the contributions to the soft $m_{H_u}^2$ and $m_{H_d}^2$ masses and/or to the μ and B_μ terms by an appropriate choice of the Θ, ξ and y parameters. For example taking $\Theta = 0$ or $\xi = 1$ leads to $\mu = B_\mu = 0$ whereas fixing $y = 1$ shuts off the extra contributions to the $m_{H_u}^2$ and $m_{H_d}^2$ soft masses. In both cases non-zero A_u and A_d are generated at the messenger scale.

With these prerequisites we are now ready to write down the full UV completion of the low-energy effective theory presented in section 2. Our idea is to simply take three copies of the second line in eq.(14):

$$W \supset X\Phi\bar{\Phi} + \lambda N H_d H_u + \frac{\kappa}{3} N^3 + \tag{25}$$

$$+ \ \lambda_{u,T} H_u D T + \lambda_{d,T} H_d \bar{D}\bar{T} + X_D D\bar{D} + \frac{X_T}{2} (T \ \bar{T}) \begin{pmatrix} a_T & a_{T\bar{T}} \\ a_{T\bar{T}} & a_{\bar{T}} \end{pmatrix} \begin{pmatrix} T \\ \bar{T} \end{pmatrix} +$$

$$+ \ \lambda_{u,P} H_u F P + \lambda_{d,P} H_d \bar{F}\bar{P} + X_F F\bar{F} + \frac{X_P}{2} (P \ \bar{P}) \begin{pmatrix} a_P & a_{P\bar{P}} \\ a_{P\bar{P}} & a_{\bar{P}} \end{pmatrix} \begin{pmatrix} P \\ \bar{P} \end{pmatrix} +$$

$$+ \ \lambda_{u,S} H_u E S + \lambda_{d,S} H_d \bar{E}\bar{S} + X_E E\bar{E} + \frac{X_S}{2} (S \ \bar{S}) \begin{pmatrix} a_S & a_{S\bar{S}} \\ a_{S\bar{S}} & a_{\bar{S}} \end{pmatrix} \begin{pmatrix} S \\ \bar{S} \end{pmatrix}$$

The four additional $SU(2)$ doublets F, E, $\bar{F}$, $\bar{E}$ have the same quantum numbers as D and $\bar{D}$ respectively. The other extra superfields $P, \bar{P}, S, \bar{S}$ are total singlets with respect to the SM gauge group. It is implicitly understood that all three pairs $D, \bar{D}$, $F, \bar{F}$ and $E, \bar{E}$ lie in $\mathbf{5} \oplus \bar{\mathbf{5}}$ copies of $SU(5)$. The extra spurions X_F and X_P acquire VEVs in their scalar and F-term components:

$$\begin{align}
X_F &= M_F(1 + \Lambda_F \theta^2) & X_E = M_E(1 + \Lambda_E \theta^2) && (26) \\
X_P &= M_P(1 + \Lambda_P \theta^2) & X_S = M_S(1 + \Lambda_S \theta^2) && (27)
\end{align}$$

To complete the description of our model we fix

$$a_{T\bar{T}} = a_{S\bar{S}} = 0 \qquad\qquad \Lambda_P/\Lambda_F = \Lambda_S/\Lambda_E = 1 \; . \qquad\qquad (28)$$

With this choice of parameters it is then clear that

1. the second line in eq.(25) produces extra contributions only to the $m_{H_u}^2$ and $m_{H_d}^2$ masses and the soft trilinear terms A_{H_u}, A_{H_d} ,

2. the third line in eq.(25) generates non-zero μ, B_μ as well as non-zero A_{H_u}, A_{H_d} but does not contribute to either $m_{H_u}^2$ or $m_{H_d}^2$,

3. the last line in eq.(25) does not contribute to any of the μ, B_μ, $m_{H_u}^2$, $m_{H_d}^2$ mass terms in the Higgs sector – it only generates non-zero soft trilinear terms.

There is an additional subtlety related to the value of the B_μ term in models with lopsided gauge mediation. As pointed out in [10] B_μ transforms under $U(1)_R$ phase rotations while the other scalar mass terms in the Lagrangian do not. In particular we can use the $U(1)_R$ symmetry to forbid a B_μ term altogether. To put it differently – the overall size of B_μ must be a free parameter which should depend on the value of some appropriately chosen coupling in the high-energy model.

What we get after integrating out the messenger superfields from eq.(25) is an NMSSM model with non-minimal soft terms in the Higgs sector. We view this as an effective theory whose UV cutoff is the lowest decoupling scale among the $M, M_D, M_T, M_F, M_P, M_E$ and M_S. For the rest of this paper we assume that this is M, i.e. the RG running always commences at this scale. After choosing the input, we run all parameters down to the EW scale using one-loop RGEs. The relevant threshold effect comes from the heavy states in the Higgs sector with mass m_A. At this scale the heavy $SU(2)$ Higgs doublet H_d is integrated out. The bino, the Higgsinos and the left handed sleptons all live at the weak scale m_Z. For the rest of the sparticle spectrum we assume a common decoupling mass scale m_{SUSY}. After minimizing the scalar Higgs potential once, we iterate the procedure, each time adjusting the values of the input parameters in order to correctly reproduce the top mass as well as the mass of the Z-boson. To be exact we impose the following two constraints at the weak scale:

1. We need to have $v = \sqrt{v_d^2 + v_u^2} = 174$ GeV in order to reproduce the correct value of the Z-boson mass m_Z (where as usual $v_u = \langle H_u^0 \rangle$ and $v_d = \langle H_d^0 \rangle$ denote the vacuum expectation values of the neutral components of the two Higgs doublets).

2. The low-energy value of the top Yukawa coupling y_t should correctly reproduce the top quark mass, i.e. $m_t = y_t(M_{\mathrm{EW}})v_u$ [1]

4   Origin and composition of the effective μ term

It has been known for quite some time that the minimal GMSB extension of the NMSSM does not lead to phenomenologically viable spectra due to the small value of the induced μ term. As noted already in [32] this problem can be restated in terms of $\det \mathcal{M}^2_{\mathrm{CP-even}}$, the determinant of the CP-even squared mass matrix in the Higgs sector. Specifically one can show that large values of the effective SUSY breaking scale, which are necessary in order to satisfy the current LHC bound on the gluino, can lead to a negative $\det \mathcal{M}^2_{\mathrm{CP-even}} < 0$.

The origin of this result can be understood analytically if one examines in more detail the following minimization condition in the Higgs sector:

$$2\,\frac{\kappa^2}{\lambda^2}\,(\,\lambda^2 s^2\,) \;=\; \lambda v^2(\kappa \sin 2\beta - \lambda) - m_N^2 + A_\lambda \lambda\, v^2\,\frac{\sin 2\beta}{2s} + \kappa A_\kappa\, s \;. \tag{29}$$

Here $s = \langle N \rangle$ is the VEV of the scalar component of N. One can argue that in minimal gauge mediation all terms on the right hand side remain relatively small. This applies in particular to the m_N^2 soft mass which is zero at the messenger scale and remains small all the way down to m_{EW} due to its small β-function. As we will see this is the key ingredient that changes within the lopsided setting. For now let us see what the implications of the aforementioned observation are. We will argue that eq.(29) imposes an upper bound on the μ term in MGM models (it restricts this mass term to several GeV). For the sake of contradiction assume that $\mu_{\mathrm{eff,\,MGM}} = \lambda s$ can be of the order $\mathcal{O}(100)$ GeV. In this case the left hand side of (29) will be unacceptably large unless one imposes $\lambda \gg \kappa$. In order to see why this relation is problematic, let us use the following approximation for $\det \mathcal{M}^2_{\mathrm{CP-even}}$ which retains only the terms with the highest power of μ (see [32]):

$$\det \mathcal{M}^2_{\mathrm{CP-even}} \simeq \mathrm{const.}\ \mu^4 \left(-4k\lambda^4 + 2\kappa^3 g^2 + 2\kappa^3 g^2 \cos(4\beta) + 8\kappa^2\lambda^3 \sin(2\beta)\right) \tag{30}$$

Given that $\lambda \gg \kappa$, it is clear that the first term dominates, leading to a negative $\det \mathcal{M}^2_{\mathrm{CP-even}} < 0$. To make things worse, the small value of the κ coupling implies a nearly massless "Goldstone" mode in the low-energy spectrum (recall that in the MGM version of the NMSSM $\frac{\kappa}{3}N^3$ is the only coupling which breaks Peccei-Quinn explicitly).

In the lopsided NMSSM the aforementioned problem is no longer present because one generically expects that the values of κ and λ will be of the same order of magnitude. Accordingly, contributions from the last three *positive* terms in eq.(30) become comparable to and can cancel out the negative contribution from the first term. In this case the determinant of the Higgs mass matrix becomes positive.

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[1] Here m_t stands for the running top mass which differs from the pole mass m_t^{pole}. To order α_s the relation between the two is given by $m_t^{pole} = m_t(m_t)\left(1 + \frac{4\alpha_s}{3\pi}\right)$.

Let us now try to understand the theoretical underpinning behind the statement that the λ and κ couplings are expected to be of the same order of magnitude. To this end we shall once again take a look at the minimization condition (29). Previously we argued that all terms on the right hand side are very small compared to λs. This applies in particular to the singlet soft mass m_N^2 which is practically zero at the messenger scale and remains very small all the way down to the EW scale due to the small β-function. In the lopsided NMSSM this picture changes due to the non-typical RG running of m_N^2, an effect which is similar to the one we encountered in the slepton and the squark sectors. To make this last statement more precise, lets take a look at the one-loop RGE for the singlet soft mass m_N^2:

$$16\pi^2 \frac{d}{dt} m_N^2 = 4\lambda^2 \left(m_{H_u}^2 + m_{H_d}^2 + m_N^2 + 4A_\lambda^2 \right) + 4\kappa^2 \left(3m_N^2 + 4A_\kappa^2 \right) \tag{31}$$

Note that the dominant $m_{H_d}^2$ term generates a large negative m_N^2 along the RG trajectory. In particular eq.(29) can be satisfied without invoking $\lambda \gg \kappa$ which solves the aforementioned determinant problem. Additionally, the large m_N^2 triggers a VEV for the singlet N which translates into a sizeable contribution to the μ term. Last but not least, the moderately large value of κ implies that PQ is no longer an approximate symmetry of the Lagrangian. Hence the quasi-Goldstone mode in the Higgs sector is avoided in a completely natural manner.

5 Low-energy spectrum
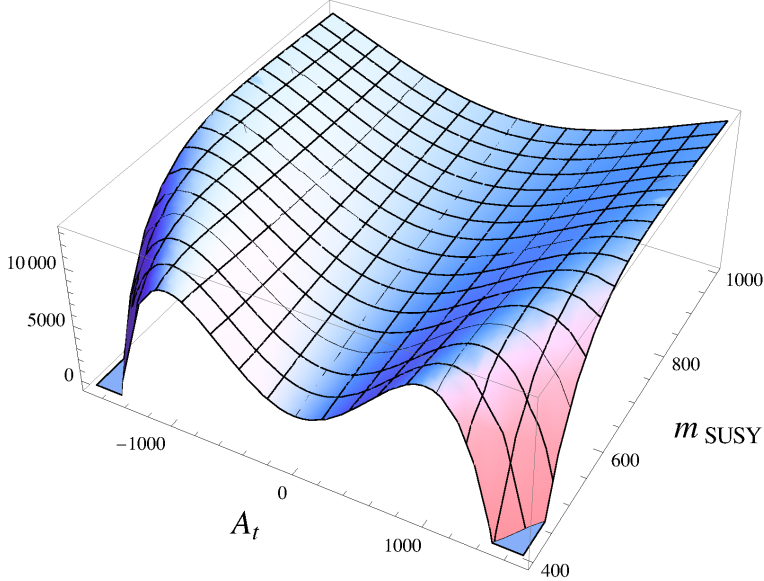
5.1 Higgs sector

The crucial issue in this section is the mass of the lightest CP-even neutral Higgs boson. Our discussion is based on the approximate one-loop formula [29]

$$\begin{aligned}
m_{h^0}^2 &= m_Z^2 \cos^2(2\beta) + \lambda^2 v^2 \sin^2(2\beta) - \frac{\lambda^2}{\kappa^2} v^2 \left(\lambda - \kappa \sin(2\beta) \right)^2 + \\
&+ \frac{3m_t^4}{4\pi^2 v^2} \left(\ln\left(\frac{m_{SUSY}^2}{m_t^2} \right) + \frac{X_t^2}{m_{SUSY}^2} \left(1 - \frac{X_t^2}{12\, m_{SUSY}^2} \right) \right) .
\end{aligned} \tag{32}$$

where $X_t = A_t - \mu \cot\beta$ is the stop mixing parameter and $m_{SUSY} \equiv \sqrt{m_{\tilde{1}}\, m_{\tilde{2}}}$ stands for the SUSY scale. Here we use eq.(32) only for illustrative purposes, the actual calculation is done using NMSSMTools [13,14]. Let us first analyze the radiative correction specified by the second line in (32). To this end we plot the magnitude of this correction as a function of A_t and m_{SUSY} in Figure 1. The symmetric ridges on the left and right hand sides denote the phenomenologically preferred region where the one-loop radiative correction is maximized.

Very recently the authors of [31] have performed a similar analysis focussing on the MSSM piece of eq.(32). It was argued that in order to lift the Higgs mass up to 125 GeV one needs a large soft trilinear term $A_t > 1$ TeV, unless of course one is prepared to

Figure 1: The radiative corrections to the Higgs mass as a function of A_t and m_{SUSY}.

tolerate superheavy stops ($\sim 5 - 10$ TeV) in the theory. This is in agreement with our plot which shows that the radiative correction is maximized for large values of the soft trilinear term. For the remainder of the paper we focus on those regions of parameter space with $|A_t| > 1$ TeV and identify $X_t \approx A_t$.

Next we discuss the NMSSM specific contributions to the tree-level $m_{h^0}^2$. Since $\tan\beta \gg 1$ we infer that $\lambda^2 v^2 \sin^2(2\beta) \approx 0$. The negative tree-level contribution due to the mixing with the singlet requires more detailed analysis. In the following we will look at several points in parameter space with $|A_t| > 1$ TeV and we will evaluate its magnitude. In order to fix the notation we introduce the three rotation angles Θ_T, Θ_P, Θ_S as well as the proportionality factors ξ_T, ξ_P, ξ_S corresponding to the last three lines in eq.(25):

$$\xi_T = 1.0 \qquad \xi_S = 1.0 \qquad \tan\Theta_P = 1.0 \tag{33}$$

We refrain from specifying $\tan\Theta_T$ and $\tan\Theta_S$ since their values will not affect any of the mass terms in the Higgs sector. Next we fix the following mass scales:

$$M_T/M_D = 1.1 \qquad M_P/M_F = 1.1 \qquad M_S/M_E = 1.1 \tag{34}$$
$$M_D = 500 \text{ TeV} \qquad M_F = 500 \text{ TeV} \qquad M_E = 500 \text{ TeV} \tag{35}$$
$$\Lambda_P = 24 \text{ TeV} \qquad \Lambda_S = 80 \text{ TeV} \tag{36}$$
$$\Lambda_F = 24 \text{ TeV} \qquad \Lambda_E = 80 \text{ TeV} \tag{37}$$

The three models (or more precisely the three points in parameter space) we consider are specified by varying the residual input parameters (Table 1).

The values of all Yukawa couplings are given at the messenger scale M. In Table 2 we evaluate the mixing effect for the three points given in Table 1.

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Table 1: Input parameters at the messenger scale M

	Point 1	Point 2	Point 3
M	10^5 GeV	10^5 GeV	10^5 GeV
ξ_P	6.75	6.75	6.75
Λ	50.0 TeV	50.0 TeV	50.0 TeV
Λ_T	63.7 TeV	63.7 TeV	63.7 TeV
Λ_D	10.4 TeV	10.4 TeV	10.4 TeV
κ	0.18	0.16	0.19
λ	0.21	0.20	0.20
$\lambda_{u,T}$	0.53	0.55	0.48
$\lambda_{d,T}$	2.29	2.24	2.51
$\lambda_{u,P}$	0.28	1.24	1.03
$\lambda_{d,P}$	0.28	1.24	1.03
$\lambda_{u,S}$	2.00	2.00	2.00
$\lambda_{d,S}$	0.00	0.00	0.00

Table 2: Mixing effect due to the presence of a gauge singlet N

	Point 1	Point 2	Point 3
$-\lambda^2/\kappa^2 \times v^2 (\lambda - \kappa \sin(2\beta))^2$	-6892 GeV^2	-7593 GeV^2	-5269 GeV^2
$m_Z^2 \cos^2(2\beta) + \lambda^2 v^2 \sin^2(2\beta)$	8190 GeV^2	8188 GeV^2	8188 GeV^2

Clearly, the mixing effect is unacceptably large for every single one of the points in parameter space we have investigated. The problem can be traced back to the fact that $\kappa < \lambda$ and $\tan\beta \gg 1$. From these we can deduce that

$$\frac{\lambda}{\kappa} v (\lambda - \kappa \sin(2\beta)) \approx \frac{\lambda^2}{\kappa} v \qquad (\lambda \gg \kappa \sin 2\beta) \tag{38}$$

i.e. the negative tree-level contribution is significant. In the following we will discuss a mechanism which will allow us to largely suppress the $\frac{\lambda^2}{\kappa^2} v^2 (\lambda - \kappa \sin(2\beta))^2$ term. The argument is based on the three minimization conditions in the Higgs sector:

$$\mu_{\text{eff}}^2 \equiv (\mu + \lambda s)^2 = -\frac{M_Z^2}{2} + \frac{m_{H_d}^2 - m_{H_u}^2 \tan^2\beta}{\tan^2\beta - 1} \tag{39}$$

$$B_{\mu,\text{eff}} \equiv (\mu + \lambda s)(A_\lambda + \kappa s) + B_\mu = \left(m_{H_u}^2 + m_{H_d}^2 + 2\lambda^2 s^2 + 2\lambda^2 v^2\right) \frac{\sin 2\beta}{2} \tag{40}$$

$$2\frac{\kappa^2}{\lambda^2}(\lambda^2 s^2) = \lambda v^2(\kappa \sin 2\beta - \lambda) - m_N^2 + A_\lambda \lambda v^2 \frac{\sin 2\beta}{2s} + \kappa A_\kappa s \tag{41}$$

Assume that we start in a region of parameter space where $\lambda > \kappa$. This would be the case if, for example, we picked one of the points listed in Table 3. In the following we will specify a procedure which will allow us to push our model into a region of parameter space where $\kappa \gg \lambda$. In this case the two terms λ and $\kappa \sin(2\beta)$ will be comparable in size and will cancel each other out leading to a highly suppressed or altogether vanishing mixing effect.

The idea of our construction is very simple – we hold the effective $\mu_{\text{eff}} = \mu + \lambda s$ term constant, while simultaneously increasing the tree-level μ piece (note that μ_{eff} is comprised of two pieces, with λs being the dynamically generated one). This can be done for instance by increasing the ξ_P parameter. Note that this procedure does not affect the $m_{H_u}^2$ and $m_{H_d}^2$ soft masses meaning that the first minimization condition (39) remains identically satisfied at every step of our construction. The third minimization condition given by eq.(41) is somewhat more problematic. Note that by increasing μ we automatically decrease the λs contribution to μ_{eff}. In particular the left hand side of eq.(41) becomes smaller in size while the right hand side remains more or less unaffected. This is easily deduced by observing that the main contribution to the r.h.s. comes from the m_N^2 soft mass term whose value is insensitive to variations of μ or λs. Thus in order to preserve the validity of eq.(41) we need to increase the κ coupling. We can then iterate this procedure until we get $\kappa \gg \lambda$.

The only open issue that remains is the validity of eq.(40). Clearly the procedure we just described affects both the left and the right hand side of this equation and there is no a priori reason why this identity should remain valid after the model parameters have been modified. Here we make use of the fact that B_μ transforms under $U(1)_R$ phase rotations while the other mass terms in the Lagrangian do not (see also [10]). In particular we can use the $U(1)_R$ symmetry to forbid a B_μ term altogether. To put it differently – the overall size of B_μ must be a free parameter which should depend on the value of some appropriately chosen coupling in the high-energy model (cf. the discussion in section 3). Hence we can always assume that B_μ is chosen in such a manner that eq.(40) is identically satisfied.

In Table 3, we list several points in parameter space with $\kappa \gg \lambda$ or $\kappa > \lambda$.

Table 3: Five points in parameter space with a large κ coupling

	Point 4	Point 5	Point 6	Point 7	Point 8
M	10^5 GeV	10^5 GeV	10^5 GeV	10^5 GeV	10^5 GeV
ξ_P	6.75	6.75	5.45	6.75	5.45
Λ	60.0 TeV	50.0 TeV	50.0 TeV	50.0 TeV	50.0 TeV
Λ_T	67.9 TeV	63.7 TeV	57.6 TeV	67.2 TeV	58.8 TeV
Λ_D	11.0 TeV	10.4 GeV	9.4 GeV	10.9 GeV	9.55 GeV
κ	0.58	0.58	0.56	0.58	0.56
λ	0.32	0.32	0.12	0.32	0.12
$\lambda_{u,T}$	0.66	0.58	0.65	0.57	0.59
$\lambda_{d,T}$	2.91	2.22	2.38	2.18	2.83
$\lambda_{u,P}$	1.20	1.30	1.45	1.27	2.05
$\lambda_{d,P}$	1.28	1.30	1.45	1.27	2.05
$\lambda_{u,S}$	2.15	2.00	2.00	1.97	1.75
$\lambda_{d,S}$	0.00	0.00	0.00	0.00	0.00

For all points in Table 3 we have chosen a large $\tan \beta = 15$. The mass of the lightest Higgs boson as well as the value and composition of the effective μ_{eff} term are given

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in Table 4. Point 4 is an example of a model in which the effective μ_{eff} term arises predominantly through the NMSSM mechanism. Conversely, for points 6 and 8 the tree-level contribution μ dominates. In fact the models represented by 6 and 8 can be viewed as the MSSM limit of the lopsided NMSSM since the effective mass term for the Higgsinos is almost entirely composed of the tree level μ term generated by the Higgs-messenger couplings.

Table 4: Higgs mass and composition of the effective μ term

	Point 4	Point 5	Point 6	Point 7	Point 8
m_{h^0}	124 GeV	124 GeV	124 GeV	124 GeV	124 GeV
μ_{eff}	314 GeV	262 GeV	140 GeV	268 GeV	262 GeV
μ	100 GeV	111 GeV	118 GeV	105 GeV	236 GeV
λs	214 GeV	151 GeV	22 GeV	163 GeV	26 GeV

5.2   Slepton and squark sector

The slepton sector of the lopsided NMSSM exhibits some peculiar features which can be traced back to the large value of the soft $m_{H_d}^2$ mass parameter. To be more precise, $m_{H_d}^2$ can have a significant impact on the running of the right and left handed soft slepton masses due to the presence of the $\frac{6}{5} Y_{\tilde{f}} g_1^2 S$ term in the β-function (see Appendix A). The aforementioned contribution is related to the induced hypercharge Fayet-Illiopoulus term in the effective action. Note that in the MSSM the impact of this term on the RG running is comparatively small. However, its significance increases dramatically once the $m_{H_d}^2$ is allowed to attain very large values. It is also important to note that $\frac{6}{5} Y_{\tilde{f}} g_1^2 S$ contributes with opposite signs to the β-functions of the right handed and left handed sleptons. Specifically, it produces a negative contribution to soft masses $m_{\tilde{L}_L}^2$ of the left-handed sleptons along the RG trajectory and a positive one to the masses $m_{\tilde{L}_R}^2$ of their right handed counterparts. As we will see shortly, this effect often leads to an atypical hierarchy $m_{\tilde{L}_R}^2 > m_{\tilde{L}_L}^2$ at the electroweak scale and can therefore reverse the ordering of the right and left handed slepton masses.

In Table 5 we list the low-energy spectrum for the same points in parameter space that we discussed at the end of section 3. The effect which reverses the ordering of the left-handed and right-handed sleptons is clearly visible both in models with low and intermediate scale gauge mediation.

In the following we will argue that the $m_{H_d}^2$ soft mass and the corresponding λ_d coupling cannot attain arbitrarily large values. One constraint arises from the requirement that the slepton masses should remain positive. Since the slepton particles are not charged under the $SU(3)$, their soft masses at the messenger scale are relatively small. Therefore the negative contribution arising from the $\frac{6}{5} Y_{\tilde{f}} g_1^2 S$ Fayet-Illiopoulus term can, at least in principle, make some of the slepton masses negative. Clearly, the tau sneutrinos are the lightest particles in the slepton sector. Imposing the condition

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Table 5: Low-energy spectrum in the slepton sector

	Point 4	Point 5	Point 6	Point 7	Point 8
$m_{\tilde{e}_R}$	499 GeV	361 GeV	371 GeV	383 GeV	414 GeV
$m_{\tilde{e}_L}$	327 GeV	307 GeV	305 GeV	297 GeV	273 GeV
$m_{\tilde{\nu}_e}$	318 GeV	297 GeV	295 GeV	287 GeV	261 GeV
$m_{\tilde{\mu}_R}$	499 GeV	361 GeV	371 GeV	383 GeV	414 GeV
$m_{\tilde{\mu}_L}$	327 GeV	307 GeV	305 GeV	297 GeV	273 GeV
$m_{\tilde{\nu}_\mu}$	318 GeV	297 GeV	295 GeV	287 GeV	261 GeV
$m_{\tilde{\tau}_1}$	260 GeV	270 GeV	263 GeV	256 GeV	205 GeV
$m_{\tilde{\tau}_2}$	419 GeV	315 GeV	308 GeV	328 GeV	334 GeV
$m_{\tilde{\nu}_\tau}$	251 GeV	267 GeV	254 GeV	249 GeV	194 GeV

Table 6: Low-energy spectrum in the squark sector

	Point 4	Point 5	Point 6	Point 7	Point 8
$m_{\tilde{u}_R}$	1552 GeV	1278 GeV	1278 GeV	1303 GeV	1264 GeV
$m_{\tilde{u}_L}$	1661 GeV	1356 GeV	1254 GeV	1385 GeV	1356 GeV
$m_{\tilde{d}_R}$	1613 GeV	1313 GeV	1310 GeV	1343 GeV	1316 GeV
$m_{\tilde{d}_L}$	1663 GeV	1358 GeV	1356 GeV	1387 GeV	1358 GeV
$m_{\tilde{s}_R}$	1613 GeV	1313 GeV	1310 GeV	1343 GeV	1316 GeV
$m_{\tilde{s}_L}$	1663 GeV	1358 GeV	1356 GeV	1387 GeV	1358 GeV
$m_{\tilde{c}_R}$	1552 GeV	1278 GeV	1278 GeV	1303 GeV	1264 GeV
$m_{\tilde{c}_L}$	1661 GeV	1356 GeV	1354 GeV	1387 GeV	1356 GeV
$m_{\tilde{t}_1}$	1033 GeV	804 GeV	777 GeV	819 GeV	814 GeV
$m_{\tilde{t}_2}$	1491 GeV	1241 GeV	1232 GeV	1254 GeV	1239 GeV
$m_{\tilde{b}_1}$	1435 GeV	1168 GeV	1158 GeV	1188 GeV	1169 GeV
$m_{\tilde{b}_2}$	1541 GeV	1266 GeV	1263 GeV	1292 GeV	1254 GeV

$m_{L_3}^2 > 0$ at the weak scale is sufficient to guarantee positivity of $m_{\tilde{\nu}_\tau}$ (and, a fortiori, of all slepton masses). For the models under consideration this does not introduce an upper bound on λ_d – even for non-perturbative values of λ_d at the messenger scale, the $m_{L_3}^2$ soft mass term remains positive at the weak scale. An actual bound on λ_d can be obtained by requiring that this coupling remains perturbative up to the GUT scale. From the one-loop RG equation for λ_d

$$\frac{d\lambda_d^2}{dt} = \frac{\lambda_d^2}{8\pi^2}\left(4\lambda_d^2 - \frac{3}{5}g_1^2 - 3g_2^2\right), \quad t = \log Q \tag{42}$$

one can deduce e.g. $\lambda_d \lesssim 1.0$ for the models with $M = 10^5$ GeV and $\lambda_d \lesssim 1.1$ for the models with $M = 10^8$ GeV. This is clearly incompatible with the values we have picked. As pointed out in [9] larger values for the λ_d coupling can be obtained by assuming that the SUSY breaking sector is strongly coupled.

There are two important effects which characterize the squark sector of the lopsided

NMSSM. The first is related to the hypercharge Fayet-Illiopoulus contribution to the β-functions of the left and right handed squarks. Note that the FI term contributes with opposite signs to the β-functions of m_U^2 and m_D^2. Since squarks are charged under the $SU(3)$ group, the RGEs are dominated by the gluino contribution $\propto g_3^2 |M_3|^2$ and the FI effect is not that prominent so as to reverse the ordering of the up and down squark masses. What it does, however, is to decrease the mass gap between right and left handed sparticles. The other effect is related to the large value of the A_t soft trilinear term. Looking at the tree-level stop mass matrix

$$\mathbf{m}_{\tilde{t}}^2 = \begin{pmatrix} m_{Q_3}^2 + m_t^2 - \left(\frac{1}{2} + \frac{2}{3} \sin^2 \theta_W \right) m_Z^2 \cos 2\beta & m_t \left(A_t - \mu \cot \beta \right) \\ m_t \left(A_t - \mu \cot \beta \right) & m_{U_3}^2 + m_t^2 + \frac{2}{3} \sin^2 \theta_W m_Z^2 \cos 2\beta \end{pmatrix} \quad (43)$$

it is immediately clear that the off-diagonal terms are large, leading to a significant splitting between the two mass eigenstates $m_{\tilde{t}_1}$ and $m_{\tilde{t}_2}$. For the sbottom mass eigenstates $m_{\tilde{b}_1}$ and $m_{\tilde{b}_2}$ the splitting is negligible. This is most easily seen by considering the respective mass matrix

$$\mathbf{m}_{\tilde{b}}^2 = \begin{pmatrix} m_{Q_3}^2 + m_b^2 - \left(\frac{1}{2} - \frac{1}{3} \sin^2 \theta_W \right) m_Z^2 \cos 2\beta & m_b \left(A_b - \mu \cot \beta \right) \\ m_b \left(A_b - \mu \cot \beta \right) & m_{D_3}^2 + m_b^2 - \frac{1}{3} \sin^2 \theta_W m_Z^2 \cos 2\beta \end{pmatrix} \quad (44)$$

and noting that $m_b A_b \ll m_t A_t$. The full spectra are listed in Table 6.

5.3 Charginos, neutralinos, gluino

The gauge singlet N adds one extra degree of freedom in the neutralino sector. To make this more precise let us introduce the fermionic component ψ_N of the superfield N. Since N is a gauge singlet, the ψ_N Weyl spinor is itself uncharged under the SM gauge group and we will refer to it as a singlino. The 5×5 neutralino mass matrix reads in the basis $(\tilde{B}^0, \tilde{W}_3^0, \tilde{H}_d^0, \tilde{H}_u^0, \psi_N)$ (cf. [29]):

$$\mathcal{M}_{\tilde{\chi}^0} = \begin{pmatrix} M_1 & 0 & -m_Z \sin\theta_W \cos\beta & m_Z \sin\theta_W \sin\beta & 0 \\ 0 & M_2 & m_Z \cos\theta_W \cos\beta & -m_Z \cos\theta_W \sin\beta & 0 \\ -m_Z \sin\theta_W \cos\beta & m_Z \cos\theta_W \cos\beta & 0 & -\mu_{\text{eff}} & -\lambda v_u \\ m_Z \sin\theta_W \sin\beta & -m_Z \cos\theta_W \sin\beta & -\mu_{\text{eff}} & 0 & -\lambda v_d \\ 0 & 0 & -\lambda v_u & -\lambda v_d & 2\kappa s \end{pmatrix}$$

The 4×4 mass matrix in the chargino sector is identical to its MSSM counterpart. Overall we get the spectra for the five points in parameter space shown in Table 7.

The neutralino χ_1^0 which arises from the mixing with the singlet is the lightest particle in the chargino/neutralino sector. A quick glance at the values in Table 7 reveals that the size of $m_{\chi_1^0}$ depends on the value of the effective μ term: In models with larger μ_{eff} the respective mass $m_{\chi_1^0}$ is larger and conversely – for smaller values of μ_{eff} within the range of ~ 100 GeV the respective neutralino mass is small leading to models with χ_1^0 as NLSP.

However, the lightest neutralino is not always NLSP. Depending on the input parameters the stau sneutrino may become lighter. An important effect, which determines

Table 7: Charginos and neutralinos

Sparticle	Point 4	Point 5	Point 6	Point 7	Point 8
$m_{\chi_1^0}$	167	103	157	166	112
$m_{\chi_2^0}$	228	- 141	178	206	- 146
$m_{\chi_3^0}$	- 229	143	- 224	232	210
$m_{\chi_4^0}$	399	210	226	- 232	223
$m_{\chi_5^0}$	462	394	390	391	394
$m_{\chi_1^\pm}$	205	124	200	208	129
$m_{\chi_2^\pm}$	399	394	389	390	394

the mass $m_{\tilde{\nu}_\tau}$, is the size of the induced FI term in the effective action, or in other words the size of the $m_{H_d}^2$ soft mass. The $m_{H_d}^2$ parameter affects the size of $m_{\tilde{\nu}_\tau}$ indirectly through the β-funcion of $m_{L_3}^2$, by decreasing the value of $m_{L_3}^2$ along the RG trajectory.

Overall we deduce that models with a large effective μ_{eff} term favour $m_{\tilde{\nu}_\tau}$ as the NLSP whereas in models with a small μ_{eff} there is a tension between the $m_{\tilde{\nu}_\tau}$ and $m_{\chi_1^0}$ masses. Point 8 is an example of a model with stau sneutrino NLSP, whereas Point 6 has χ_1^0 as NLSP.

6 Collider phenomenology

In theories with gauge mediated supersymmetry breaking the lightest supersymmetric particle (LSP) is the gravitino. The lopsided NMSSM makes no exception to this rule. In the lopsided NMSSM the next to lightest supersymmetric particle (NLSP) can be either the lightest neutralino χ_1^0 or the stau sneutrino $\tilde{\nu}$, depending on how large the effective μ term is. The crucial difference between the lopsided NMSSM and its MSSM counterpart is the fact that in the former case the effective μ term is not as strictly bounded from above.

After its production any supersymmetric state decays until the NLSP state is reached. The decay length of the NLSP is given by

$$ L \approx 10^{-2} \, \text{cm} \left(\frac{100 \, \text{GeV}}{m_{\text{NLSP}}} \right)^5 \left(\frac{\sqrt{F}}{100 \, \text{TeV}} \right)^4 \tag{45} $$

where F is the F-term component of the spurion superfield associated with the Goldstino direction. For the case $m_{\text{NLSP}} \sim 200$ GeV and assuming that $\Lambda \sim 50$ TeV we get

$L \sim 0.008$ cm for models with low-scale gauge mediation, $M = 10^5$ GeV

$L \sim 0.8$ m for models with intermediate-scale gauge mediation, $M = 10^8$ GeV

$L \sim 78000$ km for models with high-scale gauge mediation, $M = 10^{12}$ GeV

Hence, depending on the scale of supersymmetry breaking $\sqrt{F}$, the NLSP can be short- or long-lived. For the models considered in this chapter the NLSP always decays in

the detector. As far as the LHC is concerned the most relevant decay chains are those involving the gluino and the squarks. If the gluino is heavier than the squarks ($m_{\tilde{g}} > m_{\tilde{q}}$), which is always the case for the points in parameter space we have chosen, the two-body decay of the gluino into a squark and an antiquark $\tilde{g} \to \tilde{q}\bar{q}$ is kinematically possible. The squark then decays further into a quark and a chargino or neutralino. In the Tables 8 and 9 we have summarized the most relevant decay chains for gluinos and squarks for parameter point 4.

Table 8: Decays for the gluino, total width $\Gamma = 115.6$ GeV

$\tilde{g} \to$	jj	jj + l	jj + 2l	jj + 3l	jjjj	jjjj + l
Percentage	16.72 %	71.55 %	2.48 %	0.48 %	0.59 %	8.16 %

Table 9: Decays for the squarks, total width $\Gamma = 79.1$ GeV

$\tilde{q} \to$	j	j+l	j+2l	j+3l	jjj	jjj+l
Percentage	18.21 %	76.56 %	1.06 %	0.24 %	0.26 %	3.67 %

As of spring 2012, the strongest experimental lower bounds on the masses of superpartners come from searches for events with jets and missing transverse momentum at the ATLAS detector [49–51]. This signature typically arises in cascade decays of squarks and gluinos from strong pair production

$$pp \longrightarrow \tilde{g}\tilde{g}, \tilde{g}\tilde{q}, \tilde{q}\tilde{q} \,.$$

Limits are given in terms of MSUGRA/CMSSM parameters and within a simplified squark-gluino-neutralino model. The latter does not include gluino decays to third generation squarks which results in an enhanced branching ratio to first and second generation quarks. Furthermore, the neutralino is assumed massless, giving maximal phase space. These two simplifications lead to an enhanced exclusion range compared to more generic models, with $m_{\tilde{q}} \gtrsim 1.8$ TeV for $m_{\tilde{g}} \sim 1$ TeV and $m_{\tilde{q}} \gtrsim 1.3$ TeV for $m_{\tilde{g}} \gtrsim 2$ TeV [50]. However, these large exclusion bounds must be taken with a grain of salt if they are to be applied to any particular SUSY model. In the MSUGRA model for example, the exclusion limit for the average squark mass is reduced to $m_{\tilde{q}} > 1.3$ TeV for $m_{\tilde{g}} \sim 1$ TeV. The limits might be reduced even further for the model studied in this work, where decays to jets and missing transverse momentum are not the dominant decay mode of the strongly pair-produced superpartners. The situation is illustrated in Tables 8 and 9 for parameter point 4. Less than 20% of squarks and gluinos decay to jets and an invisible NLSP, giving us less than 5% of the $\tilde{q}\tilde{q}, \tilde{q}\tilde{g}, \tilde{g}\tilde{g}$ pairs decaying purely to jets and missing transverse momentum. This is due to the fact that in our NMSSM-like model, the heavy neutralinos and charginos which predominantly occur in squark and gluino decays, generically decay to leptons and sleptons. For the parameter point 4 this means

that roughly 3/4 of decays are to jets, one or more leptons, and missing energy. Consequently, roughly 80% of pair produced superpartners yield a final state with jets, at least one lepton, and missing energy. There is a number of relevant LHC searches for this type of signature [42, 43, 45, 47, 52]. One can get a rough estimate of the expected exclusion power. In the case of MSUGRA/CMSSM, the search for jets, one isolated lepton and missing energy gives roughly the same bounds as the 0-lepton channel for gluino masses $m_{\tilde{g}} < 1.2$ TeV, but considerably weaker constraints on $m_{\tilde{q}}$ for $m_{\tilde{g}} > 1.2$ TeV (see for example Figure 10 in [52]). However, we expect an enhancement of the exclusion from n-lepton final states in our model due to the suppression of 0-lepton decays. This requires a careful analysis of the various final states with leptons in the context of lopsided gauge mediation in the NMSSM, and is beyond the scope of this work. Furthermore, there are bounds on weak production of electroweak gauginos decaying to leptons [41].

If one takes the hints [53, 54] towards a ~ 125 GeV Higgs boson at the LHC at face value [55], there is a mild tension between the $\gamma\gamma$ excess and parameter points 6 and 8 which feature a lightest pseudoscalar A_1 with $2m_{A_1} < m_{h^0}$. In those cases, this decay mode dominates and $B(h^0 \to A_1 A_1) \approx 30\% \ldots 40\%$ with $B(A_1 \to b\bar{b}) \approx 90\%$, and thus all other branching fractions are reduced accordingly. At the same time, this marginally improves the situation for the WW and ZZ final state.

7 Conclusions

In the present paper we investigated a version of the Next to Minimal Supersymmetric Standard Model in which the soft masses in the Higgs sector obey the non-trivial hierarchy $m_{H_d}^2 \gg m_{H_u}^2$ at the weak scale. As argued in Chapter 5 this relation can be obtained by a suitable embedding of the low-energy model into a gauge-mediated scenario. Our investigation is prompted by a well-known result due to H. Murayama, A. de Gouvea and A. Friedland [32] which states that the $\mathbb{Z}_3$-invariant NMSSM is incompatible with the minimal version of gauge-mediated supersymmetry breaking.

We described three important effects which occur within the framework of our model. The first effect is related to the size of the VEV of the gauge singlet superfield N. Recall that in the MGM version of the NMSSM the small value of $s = \langle N \rangle$ is the very reason why the model fails to produce phenomenologically viable spectra. The lopsided extension of this scenario solves that problem by utilizing the non-standard running of the gauge singlet soft mass m_N^2. To make this statement more precise we note that the one-loop β-function of m_N^2 is comparatively large due to the presence of the dominant $m_{H_d}^2$ term. As a consequence the m_N^2 mass is shifted towards large negative values along its RG trajectory which eventually triggers a sizeable VEV for the scalar component of N.

The second effect is related to the mass m_{h^0} of the lightest CP-even Higgs boson. We showed that within the lopsided NMSSM one can easily generate large soft trilinear couplings at the messenger scale and therefore increase the one-loop radiative correction associated with top/stop loops, leading to a Higgs mass in the range of 125 GeV.

The third effect is the inverted mass hierarchy in the slepton sector. As we argued

this effect can be traced back to the atypical RG running of the soft masses of the left-handed and right-handed sleptons. To be more precise, the $m^2_{H_d}$ term enhances the $\frac{6}{5} Y_{\tilde{f}} g_1^2 \operatorname{Tr}[Y_{\tilde{f}} m^2_{\tilde{f}}]$ contribution to the one-loop β-function of the slepton masses leading to the relation $m^2_{\tilde{L}_R} > m^2_{\tilde{L}_L}$ at the electroweak scale. Something similar happens in the squark sector although the effect is much less prominent since the RG running of the colour charged sparticles is dominated by the gluino contribution.

As we showed in section 6.3 the UV completion of the lopsided NMSSM can look quite complicated. In particular we had to rely on three completely different sets of operators in order to induce the necessary extra contributions to the mass terms in the Higgs sector. It would be interesting to investigate whether a simpler realization of the lopsided NMSSM is possible.

Acknowledgements

A One-loop RGEs for the soft squark, slepton and Higgs masses in the NMSSM

In this appendix we list the one-loop RGEs for the soft mass terms in the NMSSM in the presence of a tree-level μ-term (see [29] and [18]). This is the low-energy model which is obtained after integrating out all messenger superfields in eq.(25) at their respective decoupling scales. It will be useful to introduce the following quantities:

$$X_t = m^2_{Q_3} + m^2_{U_3} + m^2_{H_u} + A_t^2 \tag{46}$$

$$X_b = m^2_{Q_3} + m^2_{D_3} + m^2_{H_d} + A_b^2 \tag{47}$$

$$X_\tau = m^2_{L_3} + m^2_{E_3} + m^2_{H_d} + A_\tau^2 \tag{48}$$

$$X_\lambda = m^2_{H_u} + m^2_{H_d} + m^2_N + A_\lambda^2 \tag{49}$$

$$X_\kappa = 3m^2_N + A_\kappa^2 \tag{50}$$

$$S \equiv \operatorname{Tr}[Y_{\tilde{f}} m^2_{\tilde{f}}] = m^2_{H_u} - m^2_{H_d} + \operatorname{Tr}[\mathbf{m_Q^2} - \mathbf{m_L^2} - \mathbf{2m_U^2} + \mathbf{m_D^2} + \mathbf{m_E^2}] \tag{51}$$

where the boldface $\mathbf{m}$'s are mass matrices in family space and $Y_{\tilde{f}}$ denotes the hyper-charge of the sfermion field $\tilde{f}$.

$$16\pi^2 \frac{d}{dt} m^2_{\tilde{Q}_i} = 2\delta_{i3}\, y_t^2 X_t + 2\delta_{i3}\, y_b^2 X_b - \frac{32}{3} g_3^2 |M_3|^2 - 6g_2^2 |M_2|^2 - \frac{2}{15} g_1^2 |M_1|^2 + \frac{1}{3}\, g_1^2 S$$

$$16\pi^2 \frac{d}{dt} m^2_{\tilde{u}_i} = 2\delta_{i3}\, y_t^2 X_t - \frac{32}{3} g_3^2 |M_3|^2 - \frac{32}{15} g_1^2 |M_1|^2 - \frac{4}{3}\, g_1^2 S$$

$$16\pi^2 \frac{d}{dt} m^2_{\tilde{d}_i} = 2\delta_{i3}\, y_b^2 X_b - \frac{32}{3} g_3^2 |M_3|^2 - \frac{8}{15} g_1^2 |M_1|^2 + \frac{2}{3}\, g_1^2 S$$

$$16\pi^2 \frac{d}{dt} m^2_{\tilde{L}_i} = 2\delta_{i3}\, y_\tau^2 X_\tau - 6g_2^2 |M_2|^2 - \frac{6}{5} g_1^2 |M_1|^2 - g_1^2 S$$

$$16\pi^2 \frac{d}{dt} m^2_{\tilde{e}_i} = 2\delta_{i3}\, y_\tau^2 X_\tau - \frac{24}{5} g_1^2 |M_1|^2 + 2\, g_1^2 S$$

$$16\pi^2 \frac{d}{dt} m^2_{H_u} = 2\, y_t^2 X_t + 2\, \lambda^2 X_\lambda - 6g_2^2 |M_2|^2 - 2g_1^2 |M_1|^2 + g_1^2 S$$

$$16\pi^2 \frac{d}{dt} m^2_{H_d} = 6\, y_b^2 X_b + 2\, y_\tau^2 X_\tau + 2\, \lambda^2 X_\lambda - 6g_2^2 |M_2|^2 - 2g_1^2 |M_1|^2 - g_1^2 S$$

$$16\pi^2 \frac{d}{dt} m^2_N = 2\, \lambda^2 X_\lambda + 2\, \kappa^2 X_\kappa \tag{52}$$

B   Coefficient functions for the induced mass terms in the Higgs sector

In this appendix we list the coefficient functions appearing in the formula for the induced mass terms in the Higgs sector:

$$
\begin{aligned}
P(x,y) &= \frac{x^2(1-y)^2}{(x^2-1)^3} \left(2(1-x^2) + (1+x^2)\log x^2 \right) \\
Q(x,y) &= \frac{x}{(x^2-1)^2} \left((x^2-1)(1-y) + (y-x^2)\log x^2 \right) \\
R(x,y) &= \frac{x}{(x^2-1)^3} \left((1-x^4)(1-y)^2 + [\,2x^2(1+y^2) - y(1+x^2)^2\,]\log x^2 \right) \\
S(x,y) &= \frac{1}{(x^2-1)^2} \left((x^2-1)(1-x^2 y) - x^2(1-y)\log x^2 \right)
\end{aligned}
\tag{53}
$$

All four functions have the property $|P(x,y)|, |Q(x,y)|, |R(x,y)|, |S(x,y)| \leq 1$. Additionally we have $P(x, y=1) = 0$ and $Q(x, y=1) = R(x, y=1)$.

References

[1] G. F. Giudice and R. Rattazzi, "Theories with gauge mediated supersymmetry breaking," Phys. Rept. **322** (1999) 419 [hep-ph/9801271].

[2] L. Alvarez-Gaume, M. Claudson, M. Wise, Nucl. Phys. **B 207** (1982) 96;
M. Dine and W. Fischler, Phys. Lett. **B 110** (1982) 227;
S. Dimopoulos, M. Dine, S. Raby, T. Scott Phys. Rev. Lett. **76** (1996) 3494-3497 .

[3] M. Dine and A. E. Nelson, "Dynamical supersymmetry breaking at low-energies," Phys. Rev. D **48** (1993) 1277 [hep-ph/9303230];
M. Dine, A. E. Nelson and Y. Shirman, "Low-energy dynamical supersymmetry breaking simplified," Phys. Rev. D **51** (1995) 1362 [hep-ph/9408384];
M. Dine, A. E. Nelson, Y. Nir and Y. Shirman, "New tools for low-energy dynamical supersymmetry breaking," Phys. Rev. D **53** (1996) 2658 [hep-ph/9507378].

[4] J. R. Ellis and D. V. Nanopoulos, "Flavor Changing Neutral Interactions In Broken Supersymmetric Theories," Phys. Lett. B **110** (1982) 44.

[5] R. D. Peccei and H. R. Quinn, "CP Conservation In The Presence Of Instantons," Phys. Rev. Lett. **38** (1977) 1440.

[6] R. D. Peccei and H. R. Quinn, "Constraints Imposed By CP Conservation In The Presence Of Instantons," Phys. Rev. D **16** (1977) 1791.

[7] G. F. Giudice and A. Masiero, "A natural solution to the μ-problem in supergravity theories," Phys. Lett. B **206** (1988) .

[8] G. R. Dvali, G. F. Giudice and A. Pomarol, "The μ-Problem in Theories with Gauge-Mediated Supersymmetry Breaking," Nucl. Phys. B **478** (1996) 31 [arXiv:hep-ph/9603238].

[9] C. Csaki, A. Falkowski, Y. Nomura and T. Volansky, "New Approach to the mu-Bmu Problem of Gauge-Mediated Supersymmetry Breaking," Phys. Rev. Lett. **102** (2009) 111801 [arXiv:0809.4492 [hep-ph]].

[10] A. De Simone, R. Franceschini, G. F. Giudice, D. Pappadopulo and R. Rattazzi, "Lopsided Gauge Mediation," JHEP **1105** (2011) 112 [arXiv:1103.6033 [hep-ph]].

[11] J. Ellis, F. Gunion, H. E. Haber, L. Roszkowski and F. Zwirner, "Higgs bosons in a nonminimal supersymmetric model," Phys. Rev. D. **39. 844** (1989)

[12] M. Drees, "Supersymmetric Models with Extended Higgs Sector," Int. J. Mod. Phys. A **4** (1989) 3635.

[13] U. Ellwanger, J. F. Gunion and C. Hugonie, "NMHDECAY: A Fortran code for the Higgs masses, couplings and decay widths in the NMSSM," JHEP **0502**, 066 (2005) [hep-ph/0406215].

[14] U. Ellwanger and C. Hugonie, "NMSPEC: A Fortran code for the sparticle and Higgs masses in the NMSSM with GUT scale boundary conditions," Comput. Phys. Commun. **177** (2007) 399 [hep-ph/0612134].

[15] F. Franke and H. Fraas, "Neutralinos and Higgs bosons in the next-to-minimal supersymmetric standard model," Int. J. Mod. Phys. A **12**, 479 (1997) [hep-ph/9512366].

[16] U. Ellwanger, M. Rausch de Traubenberg and C. A. Savoy, "Particle spectrum in supersymmetric models with a gauge singlet," Phys. Lett. B **315**, 331 (1993) [hep-ph/9307322].

[17] M. Bastero-Gil, C. Hugonie, S. F. King, D. P. Roy and S. Vempati, "Does LEP prefer the NMSSM?," Phys. Lett. B **489**, 359 (2000) [hep-ph/0006198].

[18] S. F. King and P. L. White, "Resolving the constrained minimal and next-to-minimal supersymmetric standard models," Phys. Rev. D **52**, 4183 (1995) [hep-ph/9505326]. *cite for β-function*

[19] S. A. Abel, S. Sarkar and P. L. White, "On the cosmological domain wall problem for the minimally extended supersymmetric standard model," Nucl. Phys. B **454**, 663 (1995) [hep-ph/9506359].

[20] C. Panagiotakopoulos and K. Tamvakis, "Stabilized NMSSM without domain walls," Phys. Lett. B **446**, 224 (1999) [hep-ph/9809475].

[21] J. P. Derendinger and C. A. Savoy "Quantum effects and SU(2)xU(1) breaking in supergravity gauge theories," Nucl. Phys. B **237** (1984)

[22] U. Ellwanger and C. Hugonie, "Yukawa induced radiative corrections to the lightest Higgs boson mass in the NMSSM," Phys. Lett. B **623**, 93 (2005) [hep-ph/0504269].

[23] G. Degrassi and P. Slavich, "On the radiative corrections to the neutral Higgs boson masses in the NMSSM," Nucl. Phys. B **825**, 119 (2010) [arXiv:0907.4682 [hep-ph]]. *radiative NMSSM specific one-loop corrections*

[24] U. Ellwanger and C. Hugonie, "Masses and couplings of the lightest Higgs bosons in the (M+1)SSM," Eur. Phys. J. C **25**, 297 (2002) [hep-ph/9909260].

[25] M. Masip, R. Munoz-Tapia and A. Pomarol, "Limits on the mass of the lightest Higgs in supersymmetric models," Phys. Rev. D **57**, R5340 (1998) [hep-ph/9801437].

[26] G. K. Yeghian, "Upper bound on the lightest Higgs mass in supersymmetric theories," hep-ph/9904488.

[27] A. Katz and B. Tweedie, "Signals of a Sneutrino (N)LSP at the LHC," Phys. Rev. D **81** (2010) 035012 [arXiv:0911.4132 [hep-ph]].

[28] Z. Kang, T. Li, T. Liu and J. M. Yang, "The Minimal Solution to the μ/B_μ Problem in Gauge Mediation," arXiv:1109.4993 [hep-ph].

[29] U. Ellwanger, C. Hugonie and A. M. Teixeira, "The Next-to-Minimal Supersymmetric Standard Model," Phys. Rept. **496** (2010) 1 [arXiv:0910.1785 [hep-ph]].

[30] Z. Komargodski and N. Seiberg, "mu and General Gauge Mediation," JHEP **0903** (2009) 072 [arXiv:0812.3900 [hep-ph]].

[31] P. Draper, P. Meade, M. Reece and D. Shih, "Implications of a 125 GeV Higgs for the MSSM and Low-Scale SUSY Breaking," arXiv:1112.3068 [hep-ph].

[32] A. de Gouvea, A. Friedland and H. Murayama, "Next-to-minimal supersymmetric standard model with the gauge mediation of supersymmetry breaking," Phys. Rev. D **57** (1998) 5676 [arXiv:hep-ph/9711264].

[33] G. F. Giudice and R. Rattazzi, "Living Dangerously with Low-Energy Supersymmetry," Nucl. Phys. B **757** (2006) 19 [arXiv:hep-ph/0606105].

[34] K. S. Babu, I. Gogoladze, M. U. Rehman and Q. Shafi, "Higgs Boson Mass, Sparticle Spectrum and Little Hierarchy Problem in Extended MSSM," Phys. Rev. D **78** (2008) 055017 [arXiv:0807.3055 [hep-ph]].

[35] J. D. Mason, "Gauge Mediation with a small mu term and light squarks," Phys. Rev. D **80** (2009) 015026 [arXiv:0904.4485 [hep-ph]].

[36] M. Dine, "Supersymmetry Breaking at Low Energies," Nucl. Phys. Proc. Suppl. **192-193** (2009) 40 [arXiv:0901.1713 [hep-ph]].

[37] S. P. Martin, "Extra vector-like matter and the lightest Higgs scalar boson mass in low-energy supersymmetry," Phys. Rev. D **81** (2010) 035004 [arXiv:0910.2732 [hep-ph]].

[38] S. P. Martin, "Raising the Higgs mass with Yukawa couplings for isotriplets in vector-like extensions of minimal supersymmetry," Phys. Rev. D **82** (2010) 055019 [arXiv:1006.4186 [hep-ph]].

[39] R. Barbieri, L. J. Hall, Y. Nomura and V. S. Rychkov, "Supersymmetry without a Light Higgs Boson," Phys. Rev. D **75** (2007) 035007 [arXiv:hep-ph/0607332].

[40] R. Barbieri, L. J. Hall, A. Y. Papaioannou, D. Pappadopulo and V. S. Rychkov, "An alternative NMSSM phenomenology with manifest perturbative unification," JHEP **0803** (2008) 005 [arXiv:0712.2903 [hep-ph]].

[41] G. Aad *et al.* [ATLAS Collaboration], "Search for supersymmetry in events with three leptons and missing transverse momentum in sqrt(s) = 7 TeV pp collisions with the ATLAS detector," arXiv:1204.5638 [hep-ex].

[42] G. Aad *et al.* [ATLAS Collaboration], "Search for supersymmetry with jets, missing transverse momentum and at least one hadronically decaying tau lepton in proton-proton collisions at sqrt(s) = 7 TeV with the ATLAS detector," arXiv:1204.3852 [hep-ex].

[43] G. Aad *et al.* [ATLAS Collaboration], "Search for events with large missing transverse momentum, jets, and at least two tau leptons in 7 TeV proton-proton collision data with the ATLAS detector," arXiv:1203.6580 [hep-ex].

[44] G. Aad *et al.* [ATLAS Collaboration], "Search for supersymmetry in pp collisions at sqrt(s) = 7 TeV in final states with missing transverse momentum and b-jets with the ATLAS detector," arXiv:1203.6193 [hep-ex].

[45] G. Aad *et al.* [ATLAS Collaboration], "Search for gluinos in events with two same-sign leptons, jets and missing transverse momentum with the ATLAS detector in pp collisions at sqrt(s) = 7 TeV," arXiv:1203.5763 [hep-ex].

[46] G. Aad *et al.* [ATLAS Collaboration], "Search for scalar bottom pair production with the ATLAS detector in pp Collisions at sqrts = 7 TeV," arXiv:1112.3832 [hep-ex].

[47] G. Aad *et al.* [ATLAS Collaboration], "Searches for supersymmetry with the ATLAS detector using final states with two leptons and missing transverse momentum in sqrts = 7 TeV proton-proton collisions," Phys. Lett. B **709**, 137 (2012) [arXiv:1110.6189 [hep-ex]].

[48] G. Aad *et al.* [ATLAS Collaboration], "Search for Massive Colored Scalars in Four-Jet Final States in sqrts=7 TeV proton-proton collisions with the ATLAS Detector," Eur. Phys. J. C **71**, 1828 (2011) [arXiv:1110.2693 [hep-ex]].

[49] G. Aad *et al.* [ATLAS Collaboration], "Search for squarks and gluinos using final states with jets and missing transverse momentum with the ATLAS detector in sqrt(s) = 7 TeV proton-proton collisions," Phys. Lett. B **710**, 67 (2012) [arXiv:1109.6572 [hep-ex]].

[50] [ATLAS Collaboration], "Search for squarks and gluinos with the ATLAS detector using final states with jets and missing transverse momentum and 4.7 $fb^{???1}$ of $\sqrt{s}$ = 7 TeV proton-proton collision data " ATLAS-CONF-2012-033 March 11, 2012

[51] [ATLAS Collaboration], "Hunt for new phenomena using large jet multiplicities and missing transverse momentum with ATLAS in L = 4.7 $fb^{???1}$ of $\sqrt{s}$ = 7 TeV proton-proton collisions " ATLAS-CONF-2012-037 Match 11, 2012

[52] [ATLAS Collaboration], "Further search for supersymmetry at $\sqrt{s}$ = 7 TeV in final states with jets, missing transverse momentum and one isolated lepton " ATLAS-CONF-2012-041 March 16, 2012

[53] The ATLAS Colloboration, Phys. Lett. B **710**, 383 (2012); Phys. Rev. Lett. **108**, 111803 (2012); Phys. Lett. B **710**, 49 (2012); Phys. Rev. Lett. **108**, 111802 (2012);

[54] The CMS Collaboration, Phys. Lett. B **710**, 91 (2012); Phys. Rev. Lett. **108**, 111804 (2012); Phys. Lett. B **710**, 284 (2012); arXiv:1202.4083[hep-ex]; CMS-PAS-HIG-12-001; CMS-PAS-HIG-12-007.

[55] M. Klute, R. Lafaye, T. Plehn, M. Rauch and D. Zerwas, "Measuring Higgs Couplings from LHC Data," arXiv:1205.2699 [hep-ph].